

Pion Excess, Nuclear Correlations, and the Interpretation of (\vec{p}, \vec{n}) Spin Transfer Experiments

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Abstract

Conventional theories of nuclear interactions predict a net increase in the distribution of virtual pions in nuclei relative to free nucleons. Analysis of data from several nuclear experiments has led to claims of evidence against such a pion excess. These conclusions are usually based on a collective theory (RPA) of the pions, which may be inadequate. The issue is the energy dependence of the nuclear response, which differs for theories with strong NN correlations from the RPA predictions. In the present paper, information about the energy dependence is extracted from sum rules, which are calculated for such a correlated, noncollective nuclear theory. The results lead to much reduced sensitivity of nuclear reactions to the correlations that are responsible for the pion excess. The primary example is (\vec{p}, \vec{n}) spin transfer, for which the expected effects are found to be smaller than the experimental uncertainties. The analysis has consequences for Deep Inelastic Scattering (DIS) experiments as well.

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1 Introduction

This paper addresses two related questions: First, whether the analysis of data from recent (\vec{p}, \vec{n}) experiments is in disagreement with conventional theories of nuclear structure. Second, whether the data from these experiments can be interpreted as evidence against the existence of pion excess in nuclei, as suggested by meson theories of nuclear forces. The analysis presented here will suggest that the answer to both questions is negative. Some consequences of this argument for the analysis of Deep Inelastic Scattering (DIS) experiments are also explored.

The background of the questions is as follows. The idea that nuclear pions could play a role in the EMC effect in DIS of electrons was investigated by several authors [1, 2, 3, 4, 5]. These pions are usually referred to as the pion *excess*, excluding the pions carried by free nucleons, which are already accounted for in the nucleon contribution to DIS. The original notion was that a large enhancement (of the EMC ratio) arises from collectivity of nuclei in ‘pionic’ modes, as supported by RPA models of nuclear matter [2, 5]. An alternative view [3, 4] is that the pion excess is connected to the strong two-nucleon correlations, particularly tensor and short range correlations, in the structure of nuclear ground states. We shall take this second view, which associates nuclear pions with conventional nuclear structure theory, and the meson theory of nuclear forces.

The size and origin of the pion excess could not be settled from the EMC effect itself, which is sensitive to other physical effects as well. Previously, it had been proposed [6] that the collective, or ‘precursor’ pionic modes could be studied in nucleon inelastic scattering at intermediate energies. This is best done by spin-transfer experiments, which separate the longitudinal and transverse response functions. A series of such experiments was carried out at LAMPF with polarized protons of about 500 MeV, first for (\vec{p}, \vec{p}') scattering [7], and later for (\vec{p}, \vec{n}) reactions [8, 9, 10]. These last are of most direct interest for us, since the nuclear response is purely isovector, allowing isolation of the pionic (longitudinal) mode.

The spin-transfer experiments are analyzed in terms of the two response functions, R_L (spin-longitudinal) and R_T (spin-transverse), which must be extracted from the scattering reaction data by some form of DWIA. (The longitudinal and transverse response functions for (e, e') are different functions from R_L and R_T , and are obtained differently from experiment.) The uncertainties inherent in this extraction are presumably reduced for the ratio R_L/R_T , which is expected to be a sensitive measure of the pionic collectivity. The predictions of RPA models of this collectivity [6, 11] give enhanced values of this ratio at moderate excitation energies of the order of the quasifree peak. The experimental values of this ratio [7, 8, 9, 10] do not show this collective enhancement. This lack of enhancement of the R_L/R_T ratio has been used as evidence for the absence of a pion excess. More properly, it is evidence against strong collectivity in the spin-longitudinal mode. However, strongly collective pionic modes have not been found in other studies of nuclear structure, either.

Other nuclear reactions have been analyzed in terms of the possible contribution of nuclear pions. Dimuon production rates on nuclear targets by 800 GeV protons [12] seem to show little change (per nucleon) from ^2H to Fe. This has been given as evidence against

(among other things) a pion excess in nuclei, although the size of the expected effect is based on comparison with RPA predictions with sizable collectivity. The nuclear pion excess has also been invoked as a possible source for extra total cross section in K^+ -nucleus scattering [13, 14]. Here the effects are expected to be small, with sufficient uncertainties in the analysis to make the evidence unclear. The connection of the nuclear pion excess to Compton scattering has been estimated for the deuteron and for nuclear matter with correlations [15]. The experimental input here is actually from the photonuclear sum rule, which appears to exhibit a pionic contribution.

The EMC effect, spin-transfer (\vec{p}, \vec{n}) data, and the dimuon production rates, have been taken as evidence that there is no, or at least, very little pion excess in nuclei. Miller [16], in a comment on Alulinichev's paper [13], argues that these experiments put a limit on the pion excess which is smaller than meson theory predicts [3, 4]. Bertsch, Frankfurt, and Strikman make a similar argument in an article, *Where are the Nuclear Pions* [17], concluding that meson theory fails at the relevant scales in these experiments. Brown, Buballa, Li and Wambach accept the evidence against conventional meson theory, but tell us in *Where the Nuclear Pions Are* [18], how a modified theory can fit the data.

The first purpose of the present paper is to reexamine the connection between the nuclear response functions, R_L and R_T , and conventional theories of nuclear structure, for momentum and energy transfers relevant to the recent (\vec{p}, \vec{n}) experiments. The point of view is that the ground states of nuclei and their excitation spectra are dominated by two-nucleon correlations induced by nucleon-nucleon interactions. Examples of explicit calculations of this type by the Argonne and Urbana groups are reviewed by Benhar and Pandharipande [19]; we shall rely on some results of these calculations [20] throughout this work. Consistent with this picture, we assume that there are no strongly collective pionic modes.

The main problem is that direct calculation of the response functions for fully correlated nuclei is considerably more complicated than for collective modes (in the RPA), and has not been carried out. Our method is to use information from sum rules for the response functions, which can be calculated (and have been, in Ref. [20]) to construct models of the response functions themselves. Using these models, we find that the contributions of the correlations to the response functions are smaller than the experimental uncertainties in the measured values: there is no disagreement between the data and the models.

Further results of our study of response functions relate to the question of the pion excess in nuclei. For the (\vec{p}, \vec{n}) data we shall see no disagreement with the predictions of the excess for correlated nuclei, for reasons closely related to those discussed for the response functions. For the EMC and dimuon production experiments, we must look at the method by which the pion contributions are conventionally calculated, based on the 'Sullivan' process [21]. Here we shall find that kinematic constraints reduce the contribution expected from the nuclear pions, because of the energy structure of the response functions. It does not seem that any of these experiments is sufficiently sensitive to the pion excess to present a disagreement with conventional theories of nuclear structure.

The structure of the paper is as follows. The response functions and the pion excess distributions are defined in Section 2. The energy-unweighted sum rules are discussed in

Section 3, and energy-weighted sum rules in Section 4. The response functions are separated into quasifree and correlation contributions; these are modeled in Section 5, using the sum rules to constrain the functions. The results are used to interpret the (\vec{p}, \vec{n}) data in Section 6, and the DIS structure functions in Section 7. Further discussion follows in the concluding section.

2 Response functions and pion excess

It is convenient to formulate the discussion in terms of response functions, which are defined as follows:

$$R_\alpha^a(q, \omega) = \sum_f | \langle f | \Gamma_\alpha^a(\mathbf{q}) | i \rangle |^2 \delta(\omega - E_f + E_i) \quad (1)$$

where i denotes the target ground state, f the excited nuclear state, E_i , E_f the energies of these states; \mathbf{q} and ω are the momentum and energy transfers, and Γ_α^a characterizes the nuclear excitation. We assume spherical symmetry (unpolarized or spin-zero target) so that R_α depends only on the magnitude $q = |\mathbf{q}|$.

For discussion of the spin-transfer reactions we require two response functions: longitudinal $R_L^a(q, \omega)$ and transverse $R_T^a(q, \omega)$; it is these two functions which are extracted (in principle) from the analysis of the (\vec{p}, \vec{n}) experiments [10]. For the range of energy and momentum transfers required, we assume the excitations are dominated by (nonrelativistic) nucleons, for which the $\Gamma_\alpha^a(\mathbf{q})$ are given by the following single-nucleon operators:

$$\Gamma_L^a(\mathbf{q}) = \sum_{k=1}^A (\boldsymbol{\sigma}_k \cdot \mathbf{q}) \tau_k^a \exp(i\mathbf{q} \cdot \mathbf{r}_k), \quad (2)$$

$$\Gamma_T^a(q) = \sqrt{\frac{1}{2}} \sum_{k=1}^A (\boldsymbol{\sigma}_k \times \mathbf{q}) \tau_k^a \exp(i\mathbf{q} \cdot \mathbf{r}_k), \quad (3)$$

where τ^a are the nucleon isospin (Pauli) operators, with $a = +, -, 0$, normalized so that $\tau^2 = \sum_a \tau^{-a} \tau^a$, and $\boldsymbol{\sigma}$ are the Pauli spin operators, nucleons labeled by k . For the (\vec{p}, \vec{n}) reaction, the charge transfer to the target gives $a = +$, with response functions $R_\alpha^+(q, \omega)$.

It will be easier to work in terms of the full isovector response functions,

$$R_\alpha(q, \omega) \equiv \sum_a R_\alpha^a(q, \omega). \quad (4)$$

For the case of an isoscalar ($T = 0$) target, we may write

$$R_\alpha(q, \omega) = 3R_\alpha^+(q, \omega). \quad (5)$$

The pion distribution in the nuclear target (A) is defined [3, 4] as the expectation of the number operator for pions of momentum \mathbf{q} and charge a :

$$n_A(q) = \langle i | \sum_a a_a^+(\mathbf{q}) a_a(\mathbf{q}) | i \rangle \quad (6)$$

normalized such that the total number of pions n_A is

$$n_A = \int \frac{d^3q}{(2\pi)^3} n_A(q). \quad (7)$$

The *excess* distribution is defined by the difference

$$\delta n_A(q) = n_A(q) - A n_N(q) \quad (8)$$

where $n_N(q)$ is the pion distribution for a single nucleon (isoscalar average: $(n_p + n_n)/2$).

These quantities are theory dependent, and are therefore not directly measurable, since the virtual pion field and its interactions are not uniquely defined. However, it makes sense to follow the standard pion-theoretic definitions, for comparisons with the literature, since it is in these terms that the discussion of the nuclear pion distribution have been carried out. Assuming the standard form of pseudovector coupling of pions to nucleons, the pion distribution (6) may be related to the longitudinal response function R_L [3, 22] by

$$n_A(q) = \frac{f^2 F^2(q)}{2\varepsilon_q m^2} \int_0^\infty d\omega \frac{R_L(q, \omega)}{[\varepsilon_q + \omega]^2}. \quad (9)$$

with pseudovector coupling constant f , form factor $F(q)$, pion energy $\varepsilon_q = \sqrt{q^2 + m^2}$, with m the pion mass. The excess distribution is similarly obtained

$$\delta n_A(q) = \frac{f^2 F^2(q)}{2\varepsilon_q m^2} \int_0^\infty d\omega \frac{\delta R_L(q, \omega)}{[\varepsilon_q + \omega]^2}. \quad (10)$$

with $\delta R_L = R_L^{(A)} - A R_L^{(N)}$, where $R_L^{(N)}$ is the response function for the *nucleon* (isoscalar average). One often evaluates the excess distribution in the static limit, ignoring the nuclear excitation ω compared to ε_q , in the denominator of the integral of (10). Taking $[\varepsilon_q + \omega] \simeq \varepsilon_q$, this becomes:

$$\delta n_A(q) = \frac{f^2 F^2(q)}{2\varepsilon_q^3 m^2} \int_0^\infty d\omega \delta R_L(q, \omega). \quad (11)$$

This can be evaluated in terms of a ground-state expectation value of a two-body operator, by use of a sum rule, as discussed in the following section.

In section 7 we shall discuss another, related function (see (46)) which has also been called the pion distribution function, in the context of DIS.

3 Sum rules and correlations

The strategy of this paper is to extract information about the response functions from sum rules for these quantities. This will allow us to deal with fully interacting theories of the nuclear targets, as given by the Argonne and Urbana groups [19], for which the full response functions have not been calculated.

The first sum rule relates the integrals of the response functions over ω to spin-dependent correlations in the nuclear ground state. Define the sum functions:

$$S_\alpha(q) \equiv \frac{1}{3Aq^2} \int_0^\infty d\omega R_\alpha(q, \omega) = S_\alpha^{(1)}(q) + S_\alpha^{(2)}(q). \quad (12)$$

The notation indicates that the sum separates into one-body and two-body parts, which may be expressed as expectation values of one- and two-nucleon operators; the forms follow directly from (1) - (3):

$$S_L^{(1)}(q) = \frac{1}{3Aq^2} \sum_{k=1}^A \langle i | (\boldsymbol{\sigma}_k \cdot \mathbf{q})^2 \tau_k^2 | i \rangle = 1, \quad (13)$$

$$S_L^{(2)}(q) = \frac{1}{3Aq^2} \sum_{k \neq j}^A \langle i | (\boldsymbol{\sigma}_k \cdot \mathbf{q})(\boldsymbol{\sigma}_j \cdot \mathbf{q}) \tau_k \cdot \tau_j \exp[i\mathbf{q} \cdot (\mathbf{r}_k - \mathbf{r}_j)] | i \rangle. \quad (14)$$

$$S_T^{(1)}(q) = \frac{1}{6Aq^2} \sum_{k=1}^A \langle i | (\boldsymbol{\sigma}_k \times \mathbf{q})^2 \tau_k^2 | i \rangle = 1, \quad (15)$$

$$S_T^{(2)}(q) = \frac{1}{6Aq^2} \sum_{k \neq j}^A \langle i | (\boldsymbol{\sigma}_k \times \mathbf{q})(\boldsymbol{\sigma}_j \times \mathbf{q}) \tau_k \cdot \tau_j \exp[i\mathbf{q} \cdot (\mathbf{r}_k - \mathbf{r}_j)] | i \rangle. \quad (16)$$

Integration over the nucleon positions in (14) and (16) is implied. The normalizations are chosen for convenience in comparison with the literature.

The sum functions $S_L(q)$, $S_T(q)$ have been calculated for some nuclear ground states [20] based on realistic nucleon-nucleon interactions and variationally determined correlations. Table I contains values of these functions for ^{16}O , for a range of $q < 800 \text{ MeV}/c$ [20, Fig. 2]. The fact that $S_L(q) > 1$ for $q > 300 \text{ MeV}/c$, while $S_T(q) < 1$ for all q is a direct reflection of the strong tensor correlations of nucleon pairs in the ground state.

It is useful to compare the summed functions (12) calculated for a fully interacting nucleus to the same functions for a nuclear ground state in the shell model approximation, where the only correlations come from antisymmetry. Then, the resulting functions $S^{SM}(q) < 1$. For an L-S closed-shell nucleus, as ^{16}O , $S_L^{SM}(q) = S_T^{SM}(q) = S^{SM}(q)$. Values of this function have been calculated [14, Fig. 2], [23], and are listed in Table I. Then, define the difference functions

$$\Delta S_\alpha(q) = S_\alpha(q) - S^{SM}(q), \quad (17)$$

which measure the effects of ground-state correlations beyond those of the shell model. These are also given for ^{16}O in Table I.

The pion excess distribution can be related to the sum rule for S_L , in the static approximation. Since the nucleon contribution is subtracted from the nuclear response in (10), only the two-body term (14) contributes to (11). Then the pion excess is related to $S_L^{(2)}$, the two-nucleon part of the sum rule:

$$\delta n_A(q) = \frac{3Aq^2 f^2 F^2(q)}{2\varepsilon_q^3 m^2} S_L^{(2)}(q). \quad (18)$$

Values of the pion excess distribution for ^{16}O are calculated from (18) and the tabulated values of $S_L(q)$, using $S_L^{(2)} = S_L - 1$, and are listed in the last column of Table I. The quantity tabulated is actually $q^2 \delta n_A(q)/2\pi^2 A$, whose q -integral gives the pion excess per nucleon. One can see that the distribution decreases slowly with q , which makes the integrated quantity quite sensitive to the treatment of high q , e.g. in the form factor. In addition, the excess distribution is negative for low q and positive for higher q , giving considerable cancellation in the integrated excess. The value for ^{16}O is $\delta n_A/A \simeq 0.03$. (Wiringa obtains 0.027 for a similar calculation [24].) Note that this value corresponds to a nuclear ground state with nucleons only, as given by the correlation function (14). Theories which include explicit Δ components produce values of $\delta n_A/A \simeq 0.08 - 0.1$ [3, 4, 24]. The values of $\delta n_A(q)$ are consequently more positive than those of Table I. We return to this point in the concluding section.

4 Sum rules and energy dependence

In this section we combine information that can be obtained from the sum rules of the previous section, along with that from energy-weighted sum rules, to characterize the functional dependence of the response functions.

The starting point is the separation of the response functions into two components:

$$R_\alpha(q, \omega) = R_\alpha^{SM}(q, \omega) + \Delta R_\alpha(q, \omega), \quad (19)$$

where R_α^{SM} represents the quasi-free (uncorrelated) contribution to R_α , and ΔR_α the contribution of the nuclear interactions. The functional form of R_α^{SM} is expected to be quasifree, with a peak in energy at $\omega \simeq q^2/2M$; we shall not require much more than this information. For L-S nuclei, R^{SM} should be independent of α . We require that R^{SM} obey the sum rule

$$S^{SM}(q) = \frac{1}{3Aq^2} \int_0^\infty d\omega R^{SM}(q, \omega). \quad (20)$$

Similarly, if

$$\Delta S_\alpha(q) = \frac{1}{3Aq^2} \int_0^\infty d\omega \Delta R_\alpha(q, \omega), \quad (21)$$

then (17) remains satisfied, and can be used to constrain the size of $\Delta R_\alpha(q, \omega)$ at fixed q if the shape of the ω -distribution is known.

Some information about the shape of the ω -distributions for R_α comes from the energy-weighted sum rule

$$W_\alpha(q) = \frac{1}{3Aq^2} \int_0^\infty d\omega \omega R_\alpha(q, \omega). \quad (22)$$

By a standard double-commutator calculation, this sum can be separated into one- and two-body parts [20]

$$W_\alpha(q) = \frac{q^2}{2M} + I_\alpha(q). \quad (23)$$

Some calculated values of I_L, I_T [20, Fig. 4] for ^{16}O are given in Table II.

Define the energy-weighted integrals of the separated parts of (19):

$$W^{SM}(q) = \frac{1}{3Aq^2} \int_0^\infty d\omega \omega R^{SM}(q, \omega), \quad (24)$$

$$\Delta W_\alpha(q) = \frac{1}{3Aq^2} \int_0^\infty d\omega \omega \Delta R_\alpha(q, \omega). \quad (25)$$

Clearly (25) obeys the relation

$$\Delta W_\alpha(q) = W_\alpha(q) - W^{SM}(q). \quad (26)$$

We want to estimate $\Delta W_\alpha(q)$, which can be used in turn to determine some characteristics of the energy distribution of $\Delta R_\alpha(q, \omega)$. Since $W_\alpha(q)$ can be calculated from the sum rule, we need an estimate of $W^{SM}(q)$. First define a mean excitation energy for a given q

$$E(q) = W^{SM}(q)/S^{SM}(q), \quad (27)$$

which we write in the form

$$E(q) = \frac{q^2}{2M} + U(q), \quad (28)$$

where $q^2/2M$ is the centroid of the quasifree energy distribution for a Fermi gas (with $q > 2k_F$), and $U(q)$ is the shift of the centroid from purely kinetic recoil. For a symmetric quasifree peak, $E(q)$ is the peak energy. Then (26) can be reexpressed, using (23),(27),(28):

$$\Delta W_\alpha(q) = \frac{q^2}{2M}(1 - S^{SM}(q)) + I_\alpha(q) - U(q)S^{SM}(q). \quad (29)$$

We can use the position of the empirical peak in energy transfer, in the (\vec{p}, \vec{n}) response function data at fixed q to estimate $E(q)$. For $240 \text{ MeV}/c < q < 380 \text{ MeV}/c$, the peak is found to be $\simeq 20 \text{ MeV}$ above the quasifree energy $q^2/2M$ [10], giving $U(q) \simeq 20 \text{ MeV}$. Substituting this value into (29) leads the estimated values of ΔW_L and ΔW_T given in Table II.

A similar estimate follows from equating $W^{SM}(q) = W^{RPA}(q)$, which may be considered to be a definition of $W^{SM}(q)$ based on the theorem of Thouless [25] for the energy weighted sum rule. The calculations for ^{16}O give $W^{RPA}(q) \simeq 70 \text{ MeV}$ for $q = 350 \text{ MeV}/c$ [26]. This gives an estimate of $U(q) \simeq 12 \text{ MeV}$, or $\Delta W_L(q) \simeq 50 \text{ MeV}$, compared to $\Delta W_L \simeq 44 \text{ MeV}$ in Table II.

Both estimates given have the feature that $\Delta W_\alpha(q)$ is a major fraction of $I_\alpha(q)$, or, equivalently, that $0 < U(q) \ll I_\alpha(q)$. This seems reasonable, since $U(q)$ represents the average effect of interactions in an uncorrelated (SM) nucleus, while $I_\alpha(q)$ carries the full effect of correlations, including tensor correlations. In fact, for the limiting case of $q = 0$, $\Delta W_\alpha(0) = W_\alpha(0) = I_\alpha(0)$, as is shown in Table II. This exact result follows from the fact that $R^{SM}/3Aq^2 \rightarrow 0$ for $q \rightarrow 0$ for ^{16}O , which is a singlet in the Wigner supermultiplet scheme. (In this case $U(q)$ is not defined, but does not influence $\Delta W_\alpha(0)$, since S^{SM} vanishes.)

5 Model estimates of ΔR_α

In this section we adopt simple functional forms for $\Delta R_\alpha(q, \omega)$, with parameters adjusted to reproduce the sum-rule quantities $\Delta S_\alpha(q)$ and $\Delta W_\alpha(q)$, whose numerical values are given in Tables I and II.

For the longitudinal function we assume that $\Delta R_L(q, \omega) \geq 0$ for all q, ω . Although response functions are positive (semi-)definite (see (1)), this doesn't follow in general for the differences ΔR_α . However, since $\Delta S_L(q) \geq 0$ for all q , the assumption is plausible. For this case it is useful to define a mean excitation energy for $\Delta R_L(q, \omega)$ for a given q , as

$$\omega_L(q) = \frac{\Delta W_L(q)}{\Delta S_L(q)}. \quad (30)$$

Using the estimates of ΔW_L (Table II) and calculations of $\Delta S_L(q)$ (Table I), we obtain values of $\omega_L(q)$ well in excess of 200 MeV. These values show that there is substantial contribution to the longitudinal response function for $\omega > 200$ MeV, as has already been shown by the Euclidean response calculations of Panharipande *et al.*[20]. This high- ω tail, which extends well above the quasifree response, R^{SM} , has important consequences, as shown in the following sections. The following models are taken to explore the consequences of various forms of the tail.

For the first model we take ΔR_L to be constant over some range of ω for fixed value of q , e.g., for $0 < \omega < 2\omega_L(q)$:

$$\Delta R_L(q, \omega) = \frac{\Delta S_L(q)}{2\omega_L(q)} \Theta(2\omega_L(q) - \omega). \quad (31)$$

This form has been chosen to satisfy the sum-rule differences (21) and (25), with $\omega_L(q)$ given by (30); there are no further parameters. This model has a symmetric distribution in ω about $\omega_L(q)$.

The second model is chosen not to be symmetric, with exponential dependence on ω for high values, and linear dependence for low values of ω :

$$\Delta R_L(q, \omega) = \frac{\Delta S_L(q)}{\beta^2} \omega \exp(-\omega/\beta). \quad (32)$$

This form satisfies (21) as it stands, and also (25) with $2\beta = \omega_L(q)$. This form has its maximum value at $\omega = \beta = \omega_L(q)/2$:

$$\max \Delta R_L(q) = (2/e) \frac{\Delta S_L(q)}{\omega_L(q)}. \quad (33)$$

The value will be $4/e$ times the constant value of (31), for any given q .

A model for ΔR_T must accomodate very small values of $\Delta S_T(q)$ of either sign (see Table I), together with positive values of $\Delta W_T(q)$ (see Table II). Now ΔR_T cannot be positive

definite everywhere, which requires more parameters than can be fit by the two sum rules alone. For a simple model with a sign change at some ω_0 , we take

$$\Delta R_T(q, \omega) = \frac{\Delta W_T(q)}{\omega_0^2} [-\Theta(\omega_0 - \omega) + \Theta(\omega - \omega_0)\Theta(2\omega_0 - \omega)]. \quad (34)$$

This form has a positive energy tail, and a negative value for $\omega < \omega_0$. It satisfies (25), but gives $\Delta S_T(q) = 0$, which is only approximately right; small adjustments to the two terms in brackets will correct the value. One might expect that $\omega_0 \simeq \omega_L(q)$.

6 Nuclear correlations and (\vec{p}, \vec{n}) data

In this section we use the estimates of the previous section to establish the sensitivity of the (\vec{p}, \vec{n}) data to the effects of correlations. For the range $200 \text{ MeV}/c \leq q \leq 400 \text{ MeV}/c$, which includes the momentum transfers of the experiments, most quantities needed for the estimates of sensitivity vary little with q , so we adopt the following 'typical' values from Tables I and II:

$$\Delta S_L = 0.15, \quad (35)$$

$$\Delta W_L = 44 \text{ MeV}, \quad (36)$$

which gives

$$\omega_L = 293 \text{ MeV}, \quad (37)$$

using (30). For the transverse terms we take

$$\Delta S_T = 0.0, \quad (38)$$

$$\Delta W_T = 39 \text{ MeV}, \quad (39)$$

which applies best for $200 \text{ MeV}/c \leq q \leq 300 \text{ MeV}/c$.

For the constant model (31), the value of ΔR_L is determined from (35) and (37):

$$\Delta R_L = 2.6 \times 10^{-4} \text{ MeV}^{-1}. \quad (40)$$

For the exponential model (32), the value of ΔR_L is bounded by (33):

$$\Delta R_L \leq 3.8 \times 10^{-4} \text{ MeV}^{-1}. \quad (41)$$

For the model of ΔR_T in (34), assuming that $\omega_0 = \omega_L$, the value for $\omega \leq \omega_L$ is given by

$$\Delta R_T = -4.5 \times 10^{-4} \text{ MeV}^{-1}. \quad (42)$$

The natural scale with which to compare these estimates or bounds on the values of ΔR_α are the magnitudes of the response functions $R_\alpha(q, \omega)$ extracted from the (\vec{p}, \vec{n}) experiments.

For the range $240 \text{ MeV}/c \leq q \leq 380 \text{ MeV}/c$, the data [10, Fig 2] peak somewhat above the quasifree recoil energy at values of

$$\max R_\alpha \simeq 1.2 - 1.7 \times 10^{-2} \text{ MeV}^{-1}, \quad (43)$$

dropping to values $\simeq 0.5 \times 10^{-2} \text{ MeV}^{-1}$ towards the ends of the range of energy losses: $\omega \leq 50 \text{ MeV}$ and $\omega \geq 150 \text{ MeV}$.

As might be expected, these experimental magnitudes are of the same order as the expected magnitude of the quasifree peak, which for a Fermi gas (nuclear matter) is given by

$$\max R^{SM} = \frac{3M}{4k_F q} \simeq 0.88 \times 10^{-2} \text{ MeV}^{-1} \quad (44)$$

for Fermi momentum $\simeq 200 \text{ MeV}/c$, and $q \simeq 400 \text{ MeV}/c$.

Our estimates of the correlation contributions to the response functions given by (40), (41), and (42) are then of the order of a few percent at the peak values, and less than 10% over the whole range of excitation energies, $\omega \leq 150 \text{ MeV}$. These contributions are *smaller* than the estimated uncertainties in the data quoted [10], including counting ($\leq 10\%$), experimental systematic (6% – 8%), and model uncertainties in extraction (20%, 10%).

The ratio R_L/R_T presumably has smaller systematic and model uncertainties, and has been used as a measure of the spin-dependent structure of the nuclear response [10, Fig. 1]. We can estimate this ratio, using (40) and (42), to be

$$\frac{R_L}{R_T} \simeq \frac{R^{SM} + \Delta R_L}{R^{SM} + \Delta R_T} \simeq 1 + \frac{\Delta R_L - \Delta R_T}{R^{SM}}, \quad (45)$$

which is of order 10% in the quasifree peak region. This is again no larger than the experimental uncertainties.

The conclusion is that the size of the effects to be expected from short-range and tensor interactions in the (\vec{p}, \vec{n}) response may not be large enough to be seen in the available experiments. This is in strong distinction to the large predicted effects of RPA models. The difference may be traced to the broad distribution of the high- ω tails of the response functions in the correlated nuclear theory. The breadth of these tails limits the size of the contributions to the response functions at any given value of ω ; the relation of size to breadth given by the sum functions. In contrast to the few- to 10% contributions to the response functions, integration over ω gives larger contributions to the summed functions S_α . For example, the ratio $S_L/S_T \simeq 1.25$ for $300 \text{ MeV}/c < q < 400 \text{ MeV}/c$ (See Table I, or [20, Fig. 3]).

7 Pion contributions to nuclear structure functions

The estimates made in Section 5 also have consequences for the analysis of deep inelastic scattering (DIS) or related processes on nuclear targets. The contribution of pions to DIS has been treated by several authors in terms of the ‘Sullivan’ process, in which virtual pions

are emitted by nucleons in the target[21, 2, 27, 5, 7, 18]. The quantity of interest in this theory is called the pion distribution function for the target, which may be written in terms of the longitudinal response function (5)[2, 27]:

$$f(y) = \frac{g^2}{16\pi^2} y \int_{(My)^2}^{\infty} dq^2 \int_0^{\omega_m} d\omega \frac{F^2(t) R_L(q, \omega)}{(t + m^2)^2}, \quad (46)$$

where $t = q^2 - \omega^2$ and $y = (q_z - \omega)/M$, with q_z the longitudinal component of \mathbf{q} , M is the nucleon mass, g the pi-nucleon pseudoscalar coupling constant (related to f of (9) by $(f/m) = (g/2M)$), and $F(t)$ the form factor. The upper limit of the ω -integral is given by the inequality

$$\omega_m \equiv q - yM \geq q_z - yM. \quad (47)$$

[Note that Rees *et al.* [7] use q_{\perp}^2 as the first integration variable, while Brown *et al.* [18] use t . However, for the present discussion, it is convenient to consider R_L at fixed q and ω , as in (46).]

The function defined in (46) is clearly related to the pion distribution defined in (9), although the two functions differ in their kinematic variables (y vs. q), in their energy denominators (see Jung and Miller [22]), but even more significantly, in the energy cutoff (47), as we shall see. The pion distribution function (46) is then integrated over the momentum variable y , weighted by the pion structure function $F^{\pi}(x/y)$ to obtain the pion contribution to the target structure function at x (in the convolution model). The interest is in the *excess* pion contribution over that for A free nucleons, since the excess would be an aspect of nuclear structure (in analogy to the excess distribution defined in (8) or (10)).

The main point of this section is the strong effect of the upper limit ω_m on the contribution of the nuclear response function to the ω -integral. The largest effect is to limit the contribution of $\Delta R_L(q, \omega)$, which has a high- ω tail extending well above ω_m . By contrast, the quasifree response R^{SM} , or the similar response in RPA, are less affected by the cutoff. To illustrate the effect, we shall ignore the ω -dependence of the t -dependent factors, $F(t)$ and $(t + m^2)^2$, by making a *static approximation*: $t \rightarrow q^2$. The remaining ω -integral is now over R_L alone. We define

$$J(q, y) = \frac{1}{3Aq^2} \int_0^{\omega_m} d\omega R_L(q, \omega), \quad (48)$$

which gives a measure of the pion contribution as a function of ω_m . Clearly,

$$J(q, y) \leq S_L(q). \quad (49)$$

For ω_m large enough to include all of $R^{SM}(q, \omega)$, we may write (48) in the form

$$J(q, y) = S^{SM}(q) + f \Delta S_L(q), \quad (50)$$

where f is the fraction of ΔS_L included in the integral. For the constant ΔR_L model (31), the fraction is given by

$$f_a = [\omega_m / 2\omega_L(q)], \quad (51)$$

while for the exponential ΔR_L model (32), the result (with $\beta = \omega_L/2$) is

$$f_b = [1 - \exp(\frac{-2\omega_m}{\omega_L})(\frac{2\omega_m}{\omega_L} + 1)]. \quad (52)$$

For example, consider $q = 400$ MeV/ c , for which $S_L(q) = 1.11$, which gives the largest pion excess. With $\omega_L = 293$ MeV as in (37), the fractions f_a , f_b , and the values of J_a , J_b are given in Table III, for several values of y . One sees that the values of $J(q, y)$ are reduced below 1.0 for $y \simeq 0.3$; this means that the effect of the *excess* pions has been eliminated from the integral by the kinematic cutoff, ω_m .

The result is that the effect of excess pions which are associated with short- range and tensor interactions may be sufficiently reduced by kinematic constraints to be inaccessible in DIS. As in the previous section, this result follows from the extent of the tails of the response functions. This feature has also been noticed by M. Ericson as an effect of adding nuclear correlations to an RPA theory[28]. A similar point has been made by Garcia-Recio, Nieves, and Oset [29], with regard to the estimates [13, 14] of the pion-excess contribution to K^+ -nucleus scattering, where again, a kinematic cutoff on ω limits the contribution of the excess pions (although in their paper, in RPA only).

8 Conclusions and discussion

We have found that estimates based on a conventional theory of nuclear structure for which the ground states and excitation spectra are dominated by two-nucleon correlations, give small effects in the response functions R_L and R_T (compared to the quasifree response), for $\omega \leq 150$ MeV. These effects are within the uncertainties reported for the data from the recent (\vec{p}, \vec{n}) experiments. Of course, there are uncertainties in our models used for estimates; however, we expect the order of the effects to be well predicted. We conclude that there is no clear disagreement between this data and this kind of nuclear structure theory.

The same nuclear theory predicts a pion excess distribution, which is given for ^{16}O in Table I, and an integrated excess per nucleon $\delta n_A/A \simeq 0.03$. These values are for the static approximation (18), and correspond to a contribution of the OPE potential to the binding of ^{16}O of $\langle v^\pi/A \rangle = -30.7$ MeV [19, Table I]. The theory includes only nucleons in the nuclear states. When Δ -components are coupled explicitly through a $\pi N\Delta$ vertex, the calculated excess per nucleon increases by about a factor of 3 [3, 4, 24] (again, in the static approximation). However, the Δ -contribution to the response functions would be expected to appear at higher excitations, $\omega > 300$ MeV, and would not contribute to the (\vec{p}, \vec{n}) data for $\omega \leq 150$ MeV.

There is no conflict between the predicted pion excess distributions based on the conventional correlated nuclear theory, and the response function data from (\vec{p}, \vec{n}) . That is because the former is an ω - integrated quantity, which includes all the contribution of ΔR_L at high ω , while the latter only samples that high- ω tail at fixed ω , and in a range for which the quasifree response dominates.

By contrast, RPA theories which predict strong pionic ('precursor') collectivity do seem to be in conflict with the data. In this case, the contribution to the pion excess will peak at values of ω in the experimental (quasifree) region. The evidence is not against the pion excess, but against the collective theory of its source. It seems likely that the choice of Landau parameter $g' \simeq 0.6-0.7$ (which gives such collectivity) is too low; an effective value of $g' \geq 0.9$ seems indicated.

Finally, the high- ω tails for the conventional correlated nuclear theory suggests changes in the interpretation of the DIS and dimuon production experiments, in terms of the pion contribution. The present work shows that the kinematical cutoff in ω given by (47) can reduce - even eliminate - the effect of excess pions from the structure functions. Again, this is in contrast to the prediction of the strongly collective theory. The point is that the pion excess as defined by the sum rule (18) requires integration of the energy transfer ω to values which may not be accessible in the DIS experiments. This feature has also been noticed by M. Ericson [28], and by Garcia-Recio *et al.*[29] in a related context.

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Table I: Sum functions calculated for ^{16}O and pion excess distribution, as functions of q

$q(\text{MeV}/c)$	S_L	S_T	S^{SM}	ΔS_L	ΔS_T	$\frac{q^2 \delta n}{2\pi^2 A}(\text{fm})$
0	0.12	0.10	0	0.12	0.10	0
100	0.38	0.33	0.21	0.17	0.12	-0.0087
200	0.73	0.64	0.58	0.15	0.06	-0.0210
300	0.97	0.79	0.82	0.15	-0.03	-0.0044
400	1.11	0.89	0.95	0.16	-0.06	0.0223
500	1.11	0.96	0.99	0.12	-0.03	0.0264
600	1.08	0.98	1.00	0.08	-0.02	0.0210
700	1.06	0.99	1.00	0.06	-0.01	0.0162
800	1.04	1.00	1.00	0.04	0.00	0.0110

Table II: Energy weighted sums calculated for ^{16}O as functions of q

$q(\text{MeV}/c)$	$q^2/2M(\text{MeV})$	$I_L(\text{MeV})$	$I_T(\text{MeV})$	$\Delta W_L(\text{MeV})$	$\Delta W_T(\text{MeV})$
0	0	34	34	34	34
100	5.3	37	37	37	37
200	21	44	41	41	38
300	48	52	47	44	39
400	85	59	53	44	38
500	133	65	58	47	40
600	192	68	62	48	42
700	260	69	64	49	44
800	341	68	65	48	45

Table III: Values of $J(q, y)$ and fraction f (see (50)) calculated for $q = 400 \text{ MeV}/c$, $\omega = 293 \text{ MeV}$, using two models of ΔR_L

y	$\omega_m(\text{MeV})$	f_a	J_a	f_b	J_b
0.0	400	0.68	1.06	0.76	1.07
0.1	306	0.52	1.03	0.62	1.04
0.2	212	0.36	1.01	0.42	1.02
0.3	118	0.20	0.98	0.19	0.98